

A Solution to the Small Phase Problem of Supersymmetry*

K.S. Babu and S.M. Barr

Bartol Research Institute
University of Delaware
Newark, DE 19716

bA-93-48

Abstract

It is a well-known problem that in supersymmetric models there are new CP-violating phases which, if unsuppressed, would give a neutron electric dipole moment 10^2 to 10^3 times the present experimental limit. Here we propose that these new phases are suppressed by CP invariance, which is broken spontaneously at a high scale and that this breaking shows up at low energies only through a universal phase of the gaugino masses. It is shown that this can well fit both ϵ and ϵ' of the neutral Kaon system. The electric dipole moments of the neutron and the electron should be not much below present limits. A model incorporating these ideas in a very economical way is presented.

*Supported in part by Department of Energy Grant #DE-FG02-91ER406267

In this letter we propose a model of CP violation that solves the small phase problem of supersymmetry and has testable low-energy consequences. The idea is that CP violation arises spontaneously at a high scale and is communicated to the low-energy world through a common phase of the gaugino-masses. All low-energy CP violation would be a consequence of this one non-vanishing phase angle. The ϵ parameter of the Kaon system would arise primarily from the phase of the gluino mass through the box diagram [1] shown in Fig. 1. As we shall see, to fit ϵ the phase of the gaugino mass must be $\gtrsim 3 \times 10^{-3}$, and the gluino should be relatively light ($M_{\tilde{g}} \lesssim 500 \text{ GeV}$). This leads to a value of $|\epsilon'/\epsilon| \simeq (1 \text{ to } 3) \times 10^{-3}$ (modulo hadronic matrix element uncertainties) which arises in the model dominantly via the gluino penguin graph of Fig. 2. Electric dipole moments (edm) [2] of the neutron (d_n) [3,4] and electron (d_e) [4] would be induced by the one-loop diagrams shown in Figs. 3 and 4, which turn out to be not far below the present experimental limits.

The problem that is solved by this idea is the tendency of the neutron edm arising from Fig. 3 to come out about a factor of 10^2 to 10^3 too large in models with low-energy supersymmetry [3,4]. In SUSY models there are new sources of CP violation in the A and B parameters, the μ parameter, and the gaugino masses. If CP is explicitly broken there is no reason, in general, why these phases should be small. If one assumes, as is natural, that these phases are of order unity and that the various as-yet-unobserved superparticles (gluino, squarks) have masses around 100 GeV , then one finds that $d_n \sim 10^{-22} \text{ e-cm}$, to be compared with present upper limit of $d_n \leq 10^{-25} \text{ e-cm}$.

One solution to this well-known difficulty is to assume that CP is a spon-

taneously broken symmetry. Then CP-violating parameters are finite, calculable, and, if they arise radiatively, naturally small. This general approach to the problem, which is not new, raises two issues. The first is that spontaneous CP violation leads to cosmic domain walls. These can be rendered harmless if they are “inflated away”. This requires that CP be broken at scales larger than the reheating temperature, which argues for the scale of spontaneous CP violation to be much higher than M_W .

The second issue is how the CP violation arising spontaneously at large scales is “fed down” to the Kaon system. Several “feeding-down” mechanisms have been proposed in the literature [5,6,7]. Those suggested in Ref. 5 and 6 were motivated by the desire to solve the θ -problem (the strong CP problem) using spontaneous CP violation and were therefore necessarily somewhat intricate. In any event, it was shown in Ref. 8 that these non-axion approaches to the θ -problem are fraught with difficulties in the context of supersymmetry. In our model the θ -problem is solved by the Peccei-Quinn mechanism [9], in particular by the KSVZ invisible axion [10], and thus our feeding-down mechanism can be much more straightforward than the proposals in Ref. 5 and 6.

The essential idea is that CP is spontaneously broken by the vacuum-expectation values (VEVs) of certain gauge-singlet scalar fields, which we will call S_i , $\langle S_i \rangle \gg M_W$. These VEVs give large complex masses to some vector-like fermions (needed anyway to realize the KSVZ invisible axion) which are non-singlet under the gauge group. These fermions, which we will denote $Q + Q^c$, do not mix with the known quarks and leptons (owing to their PQ charges). When Q and Q^c are integrated out, the masses of the

gauginos that couple to them will acquire a CP-violating phase at one-loop which is naturally of order 10^{-2} to 10^{-3} .

If grand unification is assumed and it is also assumed that the mass of Q and Q^c comes predominantly from these gauge-singlet contributions of $\langle S_i \rangle$, then the phases of the masses of the $SU(3)_C$, $SU(2)_L$ and $U(1)_Y$ gauginos will be very nearly equal. In this scenario the only significant CP violating phase in the low energy theory is this common gaugino mass-phase, and thus all low-energy CP-violation phenomenology is controlled by one parameter.

A model will now be presented which shows how this idea can be implemented in a particularly economical way in which the same sector does the breaking of both the CP invariance and the Peccei-Quinn symmetry.

Consider a SUSY $SU(5)$ model in which, in addition to the known matter fields, there is a $\{5\} + \{\bar{5}\}$ with Peccei-Quinn charge $-1/2$. Denote these $Q + Q^c$. Coupling to these are two $SU(5)$ -singlet superfields S_1 and S_2 , both with Peccei-Quinn charge of $+1$. A third singlet, S_3 , carries PQ charge of -2 . The superpotential of this sector is given by

$$W = Q^c Q (f_1 S_1 + f_2 S_2) + (a_{11} S_1^2 + a_{12} S_1 S_2 + a_{22} S_2^2) S_3 \quad (1)$$

where by CP invariance all the parameters are real. Consider the case where $\langle Q \rangle = \langle Q^c \rangle = \langle S_3 \rangle = 0$. Then $F_{S_1} = F_{S_2} = F_Q = F_{\bar{Q}} = 0$ are automatically satisfied. The breaking of CP can arise as a result of the $F_{S_3} = 0$ equation,

$$a_{11} S_1^2 + a_{12} S_1 S_2 + a_{22} S_2^2 = 0, \quad (2)$$

which is solved for

$$\langle S_2 \rangle = k \langle S_1 \rangle, \quad k \equiv \left(-a_{12} \pm \sqrt{a_{12}^2 - 4a_{11}a_{22}} \right) / (2a_{22}). \quad (3)$$

If $a_{12}^2 - 4a_{11}a_{22} < 0$, then there will be a non-trivial relative phase between $\langle S_1 \rangle$ and $\langle S_2 \rangle$, which breaks CP invariance.

Note that Eq. (3) leaves the magnitude of $\langle S_1 \rangle$, $\langle S_2 \rangle$ undetermined. It is easy to fix them at the desired PQ scale in several ways. For example, an extra term in the superpotential, $(S_1 \overline{S_1} - M_{PQ}^2) X$, where $\overline{S_1}$ and X are singlet fields with PQ charges of -1 and 0 respectively, would lead to $\langle S_1 \rangle \langle \overline{S_1} \rangle = M_{PQ}^2$ from the $F_X = 0$ equation. The soft SUSY term $m^2(|S_1|^2 + |\overline{S_1}|^2)$ will minimize the potential for $\langle S_1 \rangle = \langle \overline{S_1} \rangle = M_{PQ}$. For a_{11} , a_{22} and a_{12} all of the same order, one will have $\langle S_2 \rangle$ being also of order M_{PQ} . Hence the same fields that break $U(1)_{PQ}$, namely S_1 and S_2 , break CP spontaneously. The relative phase, $[arg \langle S_1 \rangle - arg \langle S_2 \rangle]$, is the source of all CP violation in the model, while the phase $[arg \langle S_1 \rangle + arg \langle S_2 \rangle]$ is essentially the invisible axion.

The “quarks” $Q + Q^c$ that implement the KSVZ axion idea are the means of feeding CP violation to the observable low-energy world. The feeding down occurs through the (s)quark-loop contribution to the gaugino mass shown in Fig. 5. There will be a soft-SUSY breaking term for the squarks of the form

$$V_{\text{soft}} = A_1 f_1 (Q^c Q S_1) + A_2 f_2 (Q^c Q S_2) + H.c. \quad (4)$$

Thus the phase appearing in the squark mass insertion in Fig. 5 is $arg (A_1 f_1 \langle S_1 \rangle^* + A_2 f_2 \langle S_2 \rangle^*)$ while that appearing in the quark mass insertion is $arg (f_1 \langle S_1 \rangle + f_2 \langle S_2 \rangle)$. If A_1 and A_2 were equal these phases would cancel and the one-loop contribution to the gaugino mass would be real. However, $A_1 \neq A_2$ in general. Even if $A_1 = A_2$ at the Planck scale (as is expected in supergravity models), they run differently if $a_{11} \neq a_{22}$ and would

be significantly different at the Peccei–Quinn scale, M_{PQ} , which we assume to be between 10^{10} and 10^{12} GeV. One finds for the phase of the gaugino mass

$$\arg(M_{\tilde{g}}) \equiv \phi = \frac{\alpha_G}{8\pi} \left(\frac{A_1 - A_2}{M_{1/2}} \right) \left\{ \frac{2f_1 f_2 |k| \sin \Delta}{f_1^2 + f_2^2 |k|^2 + 2f_1 f_2 |k| \cos \Delta} \right\}. \quad (5)$$

Here $M_{1/2}$ is the common gaugino mass at M_{GUT} , α_G the gauge coupling strength at M_{GUT} and Δ the phase of k in Eq. (3).

A number of remarks are now in order:

(i) If one neglected the effects of the running of the parameters between M_{GUT} and M_{PQ} , then the one-loop calculation of $\arg(M_{\tilde{g}})$ given in eq. (5) would be manifestly $SU(5)$ invariant, and the phases of the gluino, the Wino and the Bino would be all the same. Interestingly, and slightly non-trivially, this result remains true to one-loop order in the RGE even when the running is taken into account, as explained below.

In the momentum range $M_{PQ} \leq \mu \leq M_{GUT}$, since $SU(5)$ symmetry is not exact, the first term in eq. (1) will split into two pieces, a color-triplet (Ω) part and an $SU(2)$ -doublet (L) part:

$$W = \Omega^c \Omega (f_1 S_1 + f_2 S_2) + L^c L (f'_1 S_1 + f'_2 S_2) + \dots \quad (6)$$

Similarly, the soft SUSY breaking terms of eq. (4) will split into

$$V_{\text{Soft}} = A_1 f_1 (\Omega^c \Omega S_1) + A_2 f_2 (\Omega^c \Omega S_2) + A'_1 f'_1 (L^c L S_1) + A'_2 f'_2 (L^c L S_2) + H.c. \quad (7)$$

At and above M_{GUT} , one has $f_1 = f'_1$, $f_2 = f'_2$, and $A_1 = A'_1$, $A_2 = A'_2$. From the renormalization group equations for the various parameters of the model we find that in the momentum range between M_{GUT} and M_{PQ} ,

$$\frac{d}{dt}(A_1 - A_2) = \frac{d}{dt}(A'_1 - A'_2); \quad \frac{d}{dt} \left(\frac{f_1}{f_2} - \frac{f'_1}{f'_2} \right) \propto \left(\frac{f_1}{f_2} - \frac{f'_1}{f'_2} \right). \quad (8)$$

It follows from the above that $(A_1 - A_2) = (A'_1 - A'_2)$ and $f_1/f_2 = f'_1/f'_2$ at all scales. Combining with the scaling of gaugino masses, namely, $(\alpha_i/M_i) = (\alpha_G/M_{1/2})$, we arrive at the result that the phase of all the gauginos are identical at the PQ scale even after $SU(5)$ symmetry breaking. They will then remain to be equal down to the weak scale.

(ii) The gaugino phase, ϕ , can easily be $\gtrsim 3 \times 10^{-3}$ which is what is typically required (as will be seen below) to generate ϵ in the K meson system from the graph of Fig. 1. However, Eq. (5) shows that ϕ large enough to fit ϵ requires that the gaugino masses not be too large, a point that is important for the expected magnitudes of the neutron and electron edms. Taking $\alpha_G \simeq 1/28$ and noting that the magnitude of the function in the curly bracket of Eq. (5) is less than unity, we see that a phase angle $\phi \gtrsim 3 \times 10^{-3}$ requires $(A_1 - A_2)/M_{1/2} \gtrsim 2.1$. This is both an upper limit on the gaugino mass and a lower limit on the A parameter. Solving the RGE for the A parameters we found that $(A_1 - A_2) \lesssim 0.7A_0$, where A_0 is the universal A parameter at the Planck scale. For $A_0 = 500 \text{ GeV}$, we see that $M_{1/2} \lesssim 170 \text{ GeV}$, which after RGE corrections correspond to $M_{\tilde{g}} \lesssim 500 \text{ GeV}$ at the weak scale. The experimental lower limit on $M_{\tilde{g}} \gtrsim 150 \text{ GeV}$ implies that $A_0 \gtrsim 150 \text{ GeV}$, which could have important consequences for the electroweak symmetry breaking.

(iii) In general complex VEVs can induce phases in the A and B parameters of the ordinary sector at tree-level as noted in Ref. 6. This can be avoided if $W(\phi_i)|_{\phi_i=\langle\phi_i\rangle}$ is real. Since $\langle Q^c \rangle = \langle Q \rangle = \langle S_3 \rangle = \langle X \rangle = 0$, this condition is trivially satisfied for Eq. (1) since W evaluated at $\phi_i = \langle\phi_i\rangle$ vanishes identically.

(iv) There is no reason to expect any other phase than the common

gaugino phase to be significantly large at low energy. For example, in the KSVZ axion model such as this, where the known quarks and leptons and the Higgs superfields H_1 and H_2 have vanishing Peccei–Quinn charge, the $Q^c - Q - S_i$ sector is quite separate from the sector of ordinary matter (except through their coupling to the gauge/gaugino particles). Thus no one-loop diagram involving $\langle S_i \rangle$ contributes to μ or $B\mu$ or to the Yukawa couplings of the known quarks and leptons.

We now turn to the evaluation of the CP violating parameters ϵ , ϵ'/ϵ and the neutron and the electron electric dipole moments in the model. The $\Delta S = 2$ CP violating effective Hamiltonian is obtained from the gluino box graph of Fig. 1 (the $SU(2)$ gaugino box graph is suppressed by a factor of ~ 30 relative to the gluino box graph and thus is negligible).

$$\mathcal{H}_{\text{eff}}^{\Delta S=2} = \frac{\alpha_s^2}{10M_{sq}^2} \delta_{LR}^2 \sin 2\phi \, x f(x) \left[\frac{7}{3} \bar{s}_{R\alpha} d_L^\alpha \bar{s}_{R\beta} d_L^\beta + \frac{5}{9} \bar{s}_{R\alpha} d_L^\beta \bar{s}_{R\beta} d_L^\alpha \right] - (L \leftrightarrow R). \quad (9)$$

Here M_{sq} is the (common) squark mass, α, β are the color indices, ϕ is the phase of the gluino mass [Eq. (5)], $x = M_g^2/M_{sq}^2$ and the function $f(x)$ is defined as

$$f(x) = \frac{10}{3(1-x)^5} (9x + 9x^2 - x^3 - 6\ln x - 18x\ln x - 17) \quad (10)$$

with $f(1) = 1$. The parameter δ_{LR} is defined to be $\delta_{LR} = m_{\tilde{d}_L \tilde{s}_R}^2 / M_{sq}^2$. Since the mass-splitting among squarks is constrained phenomenologically to be small, we have treated the $\tilde{d}_L - \tilde{s}_R$ mass insertion (denoted by $m_{\tilde{d}_L \tilde{s}_R}^2$) in Fig. 1 as small perturbation. Note that other gluino graphs which do not involve $\tilde{d}_L \tilde{s}_R$ mass insertions (e.g., one with $\tilde{d}_L \tilde{s}_L$ mass insertion) are real and do not contribute to ϵ . This simplification is a consequence of the fact that only

the gluino mass has a non-vanishing phase. The contribution to $\text{Re}M_{12}$ (or Δm_K) from Fig. 1 is obtained by the interchange $\sin 2\phi \leftrightarrow \cos 2\phi$ and taking a relative plus sign between the (LR) and (RL) terms.

The ϵ parameter evaluated from Eq. (9) is given by

$$|\epsilon| = \frac{5}{54\sqrt{2}} B \eta \frac{\alpha_s^2}{M_{sq}^2} (\delta_{LR}^2 - \delta_{RL}^2) \sin 2\phi \, x f(x) \frac{f_K^2 m_K}{\Delta m_K} \left(\frac{m_K}{m_d + m_s} \right)^2, \quad (11)$$

where in the vacuum saturation method of evaluating the $K - \bar{K}$ matrix element B would be 1 by definition. η is the QCD correction factor from M_{sq} to the hadronic scale. If α_s in Eq. (11) is evaluated at the $\mu = M_{sq}$, then $\eta \simeq 1.8$ [11] for $\alpha_s \simeq 0.12$. $f_K \simeq 165 \text{ MeV}$ is the Kaon decay constant. The function $x f(x)$ is slowly varying with its value changing from 1 for $x = 1$ to 1.1 for $x = 0.1$. Using $m_s = 150 \text{ MeV}$, $m_d = 10 \text{ MeV}$ and $x = 1$, we obtain by fitting $|\epsilon| = 2.3 \times 10^{-3}$,

$$B \sin 2\phi (\delta_{LR}^2 - \delta_{RL}^2) \simeq 3.2 \times 10^{-9} \left(\frac{M_{sq}}{300 \text{ GeV}} \right)^2. \quad (12)$$

In our spontaneous-CP violation mechanism, the phase angle ϕ is naturally of order 3×10^{-3} , so that the larger of the mass-splittings, δ_{LR} or δ_{RL} must be $\sim 10^{-3}$, assuming that they are not accidentally close in value, (i.e., assuming that one of them dominates). Demanding that the contribution from the real part of Fig. 1 not be larger than the experimental value of Δm_K , we obtain

$$B \cos 2\phi (\delta_{LR}^2 + \delta_{RL}^2) \lesssim 1.0 \times 10^{-6} \left(\frac{M_{sq}}{300 \text{ GeV}} \right)^2. \quad (13)$$

From Eq. (12) and (13), we obtain the constraint $\phi \gtrsim 3 \times 10^{-3}$.

The dominant contribution to the $\Delta S = 1$ CP violating effective Hamiltonian arises from the gluino penguin graph of Fig. 2. (The $U(1)_Y$ gaugino

penguin contribution is two orders of magnitude smaller, the $SU(2)$ gauginos do not contribute directly. There are graphs involving $\widetilde{W}^+\widetilde{H}_1^-$ mixing, but these are an order of magnitude smaller.) Evaluating Fig. 2 we obtain

$$\mathcal{H}_{\text{eff}}^{\Delta S=1} = \left(\frac{7}{9}\right) \frac{1}{256\pi^2} \frac{g_3^3}{M_{sq}} (\delta_{LR} - \delta_{RL}) \sin\phi \sqrt{x} g(x) \left[\bar{s} \lambda^a i \sigma_{\mu\nu} (1 - \gamma_5) d G_{\mu\nu}^a + H.c. \right] \quad (14)$$

where

$$g(x) = \frac{2(2 + 3x - 6x^2 + x^3 + 6x \ln x)}{(1 - x)^4} \quad (15)$$

with $g(1) = 1$. We use the bag model calculation of Ref. (12) to evaluate the hadronic matrix element in Eq. (15) and obtain

$$\frac{\epsilon'}{\epsilon} = 5.6 \times 10^2 B' \eta' \sin\phi \sqrt{x} g(x) (\delta_{LR} - \delta_{RL}) \left(\frac{300 \text{ GeV}}{M_{sq}} \right) . \quad (16)$$

Here η' is the QCD correction factor, $\eta' = [\alpha_s(\mu)/\alpha_s(M)]^{0.92} \simeq 3$, where $\mu \sim 1 \text{ GeV}$ and $\alpha_s(\mu) \simeq 0.4$ has been used. B' is a factor introduced to parameterize the uncertainty in the matrix element and is defined to be 1 if the Bag model matrix elements given in Ref. (12) are exact. Combining Eq. (12) with Eq. (16), we obtain the prediction (for $x = 1$)

$$\frac{\epsilon'}{\epsilon} \simeq 2.7 \times 10^{-6} \left(\frac{B'}{B} \right) \left(\frac{M_{sq}}{300 \text{ GeV}} \right) \frac{1}{\delta_{LR} + \delta_{RL}} . \quad (17)$$

If we make the reasonable assumption that either δ_{LR} or δ_{RL} dominates the squark mass-splitting, we obtain $|\epsilon'/\epsilon| \simeq (1 \text{ to } 3) \times 10^{-3}$. This is clearly in the range suggested by experiments. Note that the sign of ϵ'/ϵ is not predicted in our model.

In supergravity models, if the minimal supersymmetric spectrum extends all the way upto the Planck scale, the squark mass-splitting will be too

small for Eq. (11) to account for ϵ . However, the MSSM spectrum is not expected to hold all the way to M_{Pl} , since the GUT threshold will in general bring in new effects [13]. A simple example is the realization of the see-saw mechanism for neutrino masses. Between the GUT scale and the Planck scale, the Dirac and Majorana neutrino matrices, with their elements not necessarily small, will contribute to the evolution of the squark mass matrix. The running in this short momentum range can result in relatively large values of the mass-splitting. A typical diagram which can generate $\tilde{d}_L \tilde{s}_R$ mixing via the neutrino Dirac mass matrix in $SU(5)$ is shown in Fig. 6, which can lead to $\delta_{LR} \sim (10^{-4} \text{ to } 10^{-3})$. It has also been emphasized [8] that the squark mass-degeneracy in supergravity models, in the absence of additional symmetries, will naturally be $\delta m_{sq}^2/M_{sq}^2 \sim \mathcal{O}(\alpha)$. In realistic string compactification scenarios, the squark degeneracy is indeed of this order.

One of the most interesting consequences of our fundamental hypothesis that all low-energy CP violation is the result of a common gaugino-phase is that both d_n and d_e are to be expected at a measurable level. Of course there are large hadronic uncertainties in d_n , but it is generally estimated that, with phases of order unity and sparticle masses of order 100 GeV, d_n from Fig. 3 will be about 10^2 to 10^3 times the experimental bound, as noted earlier. Since we require our gluino phase to be $\gtrsim 3 \times 10^{-3}$ to fit ϵ , it is natural to expect d_n to lie not far below the present bound.

Again, there are too many presently unknown SUSY parameters involved to allow a calculation of the electron edm. However, as noted in Ref. 14, if all the superparticles have comparable masses and the gaugino phases are all comparable, one would expect that $d_e \sim 10^{-2} d_n$. Thus d_e should lie not far

below 10^{-27} e-cm.

References

1. J.F. Donoghue, H.P. Nilles and D. Wyler, Phys. Lett. **128B**, 55 (1983); M.J. Duncan, Nucl. Phys. **B221**, 285 (1983); F. Gabbiani and A. Masiero, *ibid.*, **B322**, 235 (1989); J.S. Hagelin, S. Kelley and T. Tanaka, Preprint MIU-THP-92/60 (1993).
2. For reviews, see eg., S.M. Barr and W.A. Marciano, in *CP Violation*, (ed. C. Jarlskog, World Scientific, 1989); W. Bernreuther and M. Suzuki, Rev. Mod. Phys. **63**, 313 (1991).
3. W. Buchmuller and D. Wyler, Phys. Lett. **121B**, 321 (1982); J. Polchinski and M. Wise, *ibid.*, **125B**, 393 (1983); E. Franco and M. Mangano, *ibid.*, **135B**, 445 (1984).
4. F. del Aguila, M.B. Gavela, J.A. Grifols and A. Mendez, Phys. Lett. **126B**, 71 (1983).
5. A. Dannenberg, L. Hall and L. Randall, Nucl. Phys. **B271**, 574 (1986).
6. S.M. Barr and A. Masiero, Phys. Rev. **D38**, 366 (1988); S.M. Barr and G. Segre, *ibid.* **D48**, 302 (1993).
7. A. Pomarol, Phys. Rev. **D47**, 273 (1993); R. Garisto and G. Kane, Preprint TRI-PP-93-1 (1993).
8. M. Dine, R. Leigh and A. Kagan, Preprint SLAC-PUB-6090 (1993).

9. R.D. Peccei and H.R. Quinn, Phys. Rev. Lett. **38**, 1440 (1977).
10. J.E. Kim, Phys. Rev. Lett. **43**, 103 (1979); M.A. Shifman, A.I. Vainshtein and V.I. Zakharov, Nucl. Phys. **B166**, 199 (1981).
11. G. Ecker and W. Grimus, Z. Phys. **C30**, 293 (1986).
12. J.F. Donoghue, E. Golowich, B.R. Holstein and W.A. Ponce, Phys. Rev. **D23**, 1213 (1981).
13. L.J. Hall, V.A. Kostelecky and S. Raby, Nucl. Phys. **B267**, 415 (1986);
F. Gabbiani and A. Masiero, Ref. 1.
14. W. Bernreuther and M. Suzuki, Ref. 2.

Figure Captions

- Fig. 1. A box diagram whereby the phase of the gluino mass contributes to $Im(M_{12})$ in the neutral Kaon system.
- Fig. 2. The gluino penguin graph contribution to $|\epsilon'/\epsilon|$.
- Fig. 3. A contribution to the edm of the d -quark coming from the phase of the gluino mass. A similar diagram exists for the u -quark. These in turn induce an edm of the neutron of comparable magnitude.
- Fig. 4. A contribution to the edm of the electron arising from the phase of the photino mass. There are several other diagrams involving neutralinos and charginos.

- Fig. 5. The diagram by which the gaugino masses acquire a phase of order 3×10^{-3} .
- Fig. 6. A diagram contributing to the $\tilde{d}_L \tilde{s}_R$ squark mass insertion proportional to the neutrino Dirac mass matrix.

This figure "fig1-1.png" is available in "png" format from:

<http://arXiv.org/ps/hep-ph/9309249v1>

This figure "fig1-2.png" is available in "png" format from:

<http://arXiv.org/ps/hep-ph/9309249v1>

This figure "fig1-3.png" is available in "png" format from:

<http://arXiv.org/ps/hep-ph/9309249v1>

This figure "fig1-4.png" is available in "png" format from:

<http://arXiv.org/ps/hep-ph/9309249v1>

This figure "fig1-5.png" is available in "png" format from:

<http://arXiv.org/ps/hep-ph/9309249v1>